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On M-9-branes

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Abstract. We discuss some properties of the conjectured M-9-brane. We investigate both the worldvolume action as well as the target space solution. The worldvolume action is given by a gauged sigma model which, via dimensional reduction and duality, is shown to be related to the worldvolume actions of the branes of ten-dimensional superstring theory. The effective tension of the M-9-brane scales as $(R_{11})^3$, where R_{11} is the radius of an S^1 -isometry direction. This isometry enables us to add a cosmological constant to eleven-dimensional supergravity.

The target space solution corresponding to the M-9-brane is a (wrapped) domain wall solution of massive eleven-dimensional supergravity. This solution breaks half of the bulk supersymmetry. We consider both single M-9-branes as well as a system of two M-9-branes. In both cases one can define regions in spacetime, separated by the domain walls, with zero cosmological constant. In these regions the limit $R_{11} \rightarrow \infty$ can be taken in which case the M-9-brane is unwrapped and a massless eleven-dimensional supergravity theory is obtained.

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1. Introduction

Given the fact that 11-dimensional supergravity arises in the strong-coupling limit of type IIA superstring theory [1], one would expect that all type IIA branes arise as solutions of the equations of motion corresponding to 11-dimensional supergravity [2]. This is indeed the case for all known branes of type IIA superstring theory except for the D-8-brane which arises as a domain wall solution in ten dimensions [3, 4] and whose 11-dimensional origin so far has remained unclear (see figure 1). It has been conjectured that the 11-dimensional origin of the D-8-brane is an M-9-brane[†]. In this work we investigate some properties of this conjectured M-9-brane.

One indication that indeed an M-9-brane exists in 11 dimensions comes from a study of the M-theory superalgebra [6, 7]:

$$\{Q_\alpha, Q_\beta\} = (\Gamma^M C)_{\alpha\beta} P_M + \frac{1}{2} (\Gamma_{MN} C)_{\alpha\beta} Z^{MN} + \frac{1}{5!} (\Gamma_{MNPQR})_{\alpha\beta} Y^{MNPQR}. \quad (1)$$

The time component of the 2-form central charge Z suggests a 9-brane which breaks half of the bulk supersymmetry according to (η is the bulk supersymmetry parameter)

$$(1 - \Gamma_{0123456789})\eta = 0. \quad (2)$$

From a 10-dimensional point of view (reducing over x^{10}) this is exactly the chirality condition

$$(1 - \Gamma_\star)\eta = 0, \quad (3)$$

[†] For earlier discussions on the conjectured M-9-brane, see [4, 5].

with $\Gamma_\star = \Gamma_{10} = \Gamma_0 \Gamma_1 \cdots \Gamma_9$. Hence the worldvolume theory on the 9-brane has $N = 1$ chiral supersymmetry. This 10-dimensional worldvolume theory occurs in the Hořava-Witten description of the $E_8 \times E_8$ heterotic string [8].

It is nontrivial to find a solution of 11-dimensional supergravity that describes the expected M-9-brane. The double dimensional reduction of such an M-9-brane must lead to the D-8-brane. However, the D-8-brane solution arises in massive IIA supergravity [9] which is not directly obtainable by reduction of $d = 11$ supergravity. The reason for this is that massive IIA supergravity contains a cosmological constant proportional to m^2 where m is a parameter with the dimension of a mass[†]:

$$S_{\text{massive IIA}}(m) \sim \int d^{10}x \sqrt{|g|} \left\{ e^{-2\phi} [R - 4(\partial\phi)^2] + \frac{1}{2}m^2 + \cdots \right\}. \quad (4)$$

There does not exist a corresponding massive extension of the M-theory $d = 11$ supergravity theory (for a recent discussion of this point, see [10]). There is a simple explanation for this. The only massive extension one can write down reads

$$S_{\text{massive } d=11} \sim \int d^{11}x \sqrt{|g|} \left\{ R + \frac{1}{2}m^2 + \cdots \right\}. \quad (5)$$

However, the reduction of the cosmological constant in the above action leads to a cosmological constant in $d = 10$ whose dilaton coupling (in string-frame) does not vanish, as in (4).

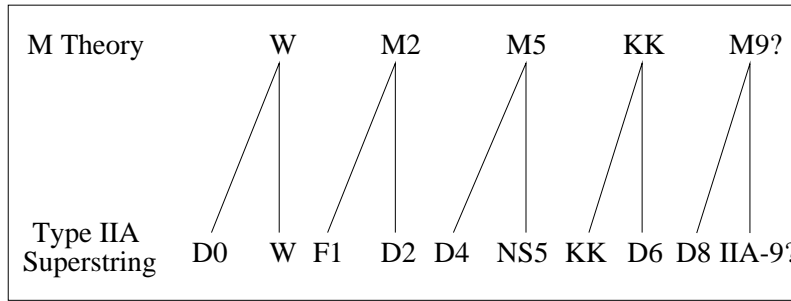


Figure 1. Relation between type IIA branes and M-theory solutions. Vertical lines imply direct dimensional reduction, diagonal lines double dimensional reduction. Besides type IIA D- p -branes ($p = 0, 2, 4, 6, 8$) and M- p -branes ($p = 2, 5$), the figure contains: waves (W); the fundamental string (F1); Kaluza–Klein monopoles (KK); and a (conjectured) IIA-9-brane. The oxidation of the D-8-brane leads to the M-9-brane discussed in the text.

One way to avoid this obstruction is to consider a $d = 11$ supergravity background with an isometry generated by a (spacelike) Killing vector k^μ [11]. This possibility was originally motivated by considering the M-theory origin of the D-2-brane in a massive background [12–14]. As is well-known, the worldvolume action of the massless D-2-brane can be rewritten as a massless M-2-brane action after a worldvolume Poincaré dualization of the Born–Infeld (BI) 1-form V into an embedding coordinate X [15]:

$$dX = *dV. \quad (6)$$

In a massive background the D-2-brane action obtains an extra Chern–Simons term proportional to m [16, 17] and, on-shell, the duality relation (6) gets replaced by

$$(dX - mV) = *dV. \quad (7)$$

[†] The parameter m can be positive or negative. The sign of the cosmological constant is determined by T-duality (see below) and is such that for constant dilaton $\phi = \phi_0$ the vacuum solution to the Einstein equations is anti-de Sitter spacetime.

It has been shown [13, 14] that in order to perform the duality transformation off-shell one must introduce an auxiliary worldvolume 1-form A and replace the Chern–Simons term by[†]

$$mVdV \rightarrow m\left(\frac{1}{2}AdA - AdV\right). \quad (8)$$

The dualization of the BI 1-form V into an embedding coordinate X then leads to a massive M-2-brane action which is gauged sigma model. The action is obtained from the massless M-2-brane action by replacing the worldvolume derivatives by covariant derivatives:

$$dX^\mu \rightarrow dX^\mu - Ak^\mu, \quad (9)$$

where A is the auxiliary 1-form introduced above and k^μ is a Killing vector. Thus the construction of the massive M-2-brane action requires the existence of a Killing vector.

In view of the above discussion it is natural to assume the existence of the same Killing vector k^μ in order to construct an 11-dimensional supergravity theory with a cosmological constant. Given this Killing vector one can define an extra scalar[‡]

$$|k|^2 = -k^\mu k^\nu g_{\mu\nu} = (R_{11})^2. \quad (10)$$

Using this scalar one can modify the Lagrangian (5) such that upon reduction it leads to a cosmological constant with the correct dilaton coupling:

$$S_{\text{massive M}} \sim \int d^{11}x \sqrt{|g|} \left\{ R + \frac{1}{2}|k|^4 m^2 + \dots \right\}. \quad (11)$$

This leads to the definition of ‘massive $d = 11$ supergravity’. The bosonic part of the action has been given in [11].

The massive supergravity theory defined above is not a proper $d = 11$ supergravity theory in the sense that it is only defined for backgrounds with a vanishing Lie derivative. *A priori* we cannot exclude the possibility that such backgrounds are special solutions of a yet to be constructed proper $d = 11$ supergravity theory with a cosmological constant. Given the no-go theorem of [10] this seems unlikely to be the case. We therefore continue our analysis using the massive $d = 11$ supergravity theory defined above. In fact, as we will show in this work, one can relax the restriction of vanishing Lie derivative to be valid only in certain regions of the $d = 11$ spacetime which are separated by domain wall M-9-brane solutions. For this to work we must use the fact that the cosmological constant is not necessarily constant everywhere but can be taken to be *piecewise constant* [3, 4]. In the case of massive IIA supergravity this is seen by using a formulation where the mass parameter m is replaced by a 9-form potential $C^{(9)}$ with curvature $G^{(10)}$ [4]§:

$$S_{\text{massive IIA}}(C^{(9)}) \sim \int d^{10}x \sqrt{|g|} \left\{ e^{-2\phi} [R - 4(\partial\phi)^2] - \frac{1}{2 \times 10!} (G^{(10)})^2 + \dots \right\}. \quad (12)$$

This 9-form potential naturally appears in the low-energy limit of type IIA superstring theory [18]. The sign of the kinetic term for $C^{(9)}$ follows, via T-duality, from the (standard) sign of the kinetic term for the other Ramond–Ramond p -form potentials $C^{(p)}$ ($p < 9$). This in turn fixes the sign of the cosmological constant in (4). More precisely, the equation of motion of $C^{(9)}$

$$d * G^{(10)} = 0 \quad (13)$$

is solved for by

$$G^{\mu_1 \dots \mu_{10}}(C) = \frac{1}{\sqrt{|g|}} \epsilon^{\mu_1 \dots \mu_{10}} C, \quad (14)$$

[†] Note that on-shell $dA = dV$ and the two formulations of the Chern–Simons term coincide.

[‡] We use a mostly minus signature $(+, -, \dots, -)$ so that $|k|^2$, as defined below, is positive for a spacelike Killing vector k^μ . We work in units in which the 11-dimensional Planck-length is equal to one.

§ In the case of massive $d = 11$ supergravity one must replace m by a 10-form potential.

where c is an integration constant. Comparing with (4) we find that $c = \pm m$.

The expression of $G^{(10)}$ given in (14) is not the most general solution of (13). In the presence of a domain wall, the integration constant c can be *piecewise constant* [4]. This possibility is excluded in the formulation without the 9-form potential where the parameter m is constant everywhere. In particular, the presence of a domain wall allows us to define the region of $d = 11$ spacetime at one side of the domain wall to have zero cosmological constant, i.e. $m = 0$. This observation will be of use later.

In this work we will use the massive $d = 11$ supergravity theory to investigate the oxidation of the D-8-brane target space solution into an M-9-brane target space solution. The organization of this paper is as follows. In section 2 we first discuss some properties of the M-9-brane worldvolume action. In particular, we discuss how the tension scales with R_{11} and how the leading Nambu–Goto term of the M-9-brane worldvolume action is related, via reduction and duality, to the worldvolume actions of the branes in 10-dimensional string theory. In section 3 we discuss some properties of the D-8-brane target space solution. In section 4 we investigate the corresponding M-9-brane target space solution. We discuss both the single M-9-brane as well as the two domain wall system. The supersymmetry properties of the D-8-brane and M-9-brane are discussed in section 5. Finally, in section 6 we present our conclusions.

2. The M-9-brane worldvolume action

One way to see that it is nontrivial to construct the worldvolume action for an M-9-brane moving in 11 *decompactified* dimensions is to consider the scaling of the tension with R_{11} . It is instructive to do this analysis for all the branes of M-theory (see figure 1). Our starting point is the gauged σ -model approach. This approach enables us to reformulate p -branes moving in a $d = 10$ target spacetime as p -branes moving in a $d = 10$ submanifold of a $d = 11$ target spacetime. In the latter case the tension of the p -brane is measured in terms of $d = 11$ quantities. The gauged σ -model plays a crucial role in the construction of the worldvolume action of the Kaluza–Klein (KK) monopole [12].

Consider the Nambu–Goto part of a general type IIA p -brane in 10 dimensions with dilaton coupling parameter α :

$$S(d = 10) = -T \int d^{p+1}x \, e^{\alpha\phi} \sqrt{|\partial_i X^\mu \partial_j X^\nu g_{\mu\nu}|}. \quad (15)$$

This action describes the dynamics of a p -brane moving in 10 dimensions. We next write down an equivalent action describing the same p -brane but now moving in a 10-dimensional subspace of an 11-dimensional manifold with an isometry generated by a spacelike Killing vector $\hat{k}^{\hat{\mu}}$ (we indicate 11-dimensional fields and indices with a hat):

$$S(d = 11) = -T \int d^{p+1}\hat{x} \, |\hat{k}|^\beta \sqrt{|\partial_i \hat{X}^{\hat{\mu}} \partial_j \hat{X}^{\hat{\nu}} \hat{\Pi}_{\hat{\mu}\hat{\nu}}|}. \quad (16)$$

The projector

$$\hat{\Pi}_{\hat{\mu}\hat{\nu}} = \hat{g}_{\hat{\mu}\hat{\nu}} + |\hat{k}|^{-2} \hat{k}_{\hat{\mu}} \hat{k}_{\hat{\nu}} \quad (17)$$

restricts the dynamics of the brane to the space orthogonal to the isometry direction. The effective tension scales as $|\hat{k}|^\beta = (R_{11})^\beta$, where β is a scaling parameter to be determined below. Note that the $d = 11$ background metric has vanishing Lie derivative with respect to the Killing vector \hat{k} .

The two actions (15) and (16) are related to each other via a direct dimensional reduction over the isometry direction. To perform this reduction it is convenient to use coordinates

$\hat{\mu} = (\mu, z)$ adapted to the isometry. In this basis the only non-vanishing components of the projector $\hat{\Pi}_{\hat{\mu}\hat{\nu}}$ are

$$\hat{\Pi}_{\mu\nu} = |\hat{k}|^{-1} g_{\mu\nu}. \quad (18)$$

Using the relation

$$\hat{g}_{zz} = -|\hat{k}|^2 = -e^{\frac{4}{3}\phi} \quad (19)$$

we find that the dilaton coupling parameter α of the 10-dimensional p -brane and the scaling parameter β of the corresponding 11-dimensional p -brane are related to each other as follows:

$$\alpha = -\frac{1}{3}(p+1) + \frac{2}{3}\beta. \quad (20)$$

For the standard branes of type IIA superstring theory the dilaton coupling parameter α takes on three different values (the discussion below is summarized in figure 2)[†]:

- (a) $\alpha = 0$ for fundamental objects like the F1 string.
- (b) $\alpha = -1$ for the D- p -branes ($p = 0, 2, 4, 6, 8$).
- (c) $\alpha = -2$ for solitonic objects like the NS5 brane.

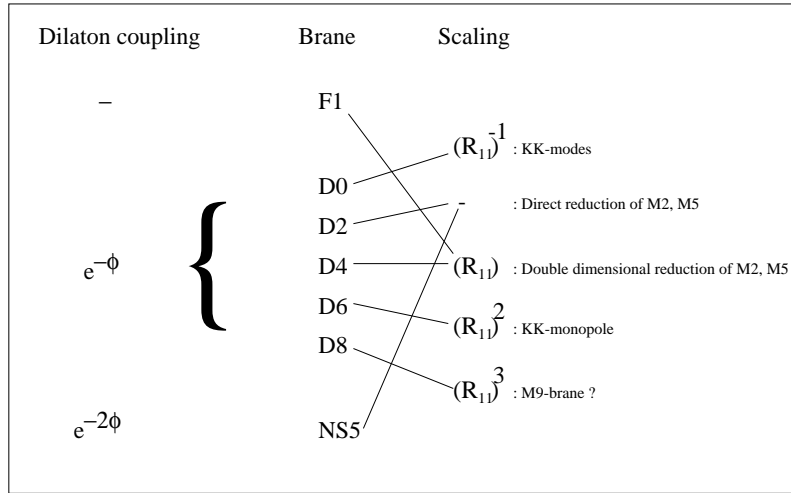


Figure 2. The scaling of type IIA branes. The type IIA-branes are divided into three classes according to their dilaton coupling. The lines in the figure indicate for each brane how the corresponding M-brane tension scales with R_{11} .

This leads to the following five different values of the scaling parameter β :

- (a) $\beta = -1$. This is the appropriate scaling for the KK modes coming from $d = 11$. In $d = 10$ they give rise to the D-0-branes.
- (b) $\beta = 0$. In this case we are dealing with the D-2-brane and NS5-brane which are related via a direct dimensional reduction to the M-2-brane and M-5-brane, respectively. In these cases the tension does not scale with R_{11} and no extra worldvolume direction is developed[‡].

[†] We do not consider here purely gravitational branes like waves and monopoles. The KK monopole will be discussed below.

[‡] In order to rewrite the D-2-brane (NS5-brane) action as an M-2-brane (M-5-brane) action one must dualize the BI 1-form into an embedding coordinate (D-2-brane) or use the extra worldvolume scalar (NS5-brane).

- (c) $\beta = 1$. This is what we expect from the F1-string and D-4-brane since they follow via a double dimensional reduction from the M-2-brane and M-5-brane, respectively. The scaling of the tension with R_{11} signatures the development of an extra worldvolume direction. Absorbing the R_{11} -factor into the Nambu–Goto part, the gauged sigma model action can be rewritten as the usual worldvolume action of the M-2-brane or M-5-brane[†].
- (d) $\beta = 2$: This happens for the D-6-brane. The D-6-brane can be obtained from the KK monopole in $d = 11$ when reduced over the U(1) isometry direction in the transversal Taub–NUT space. This isometry direction cannot be interpreted as a worldvolume direction [7]. Since the monopole cannot move in the isometry direction the corresponding scalar has to be gauged away explicitly. This leads to a gauged σ -model worldvolume action for the KK monopole [12]. The scaling of the M-theory monopole tension with $(R_{11})^2$ shows that the isometry direction must be compact: the D-6-brane cannot be oxidized to a brane moving in 11 decompactified dimensions.
- (e) $\beta = 3$ [19, 20]. This case corresponds to the D-8-brane studied in this work. We see that the M-9-brane, if it exists, cannot live in a decompactified $d = 11$ manifold, but, like the M-theory monopole it needs a compact isometry direction.

Before discussing the M-9-brane worldvolume action it is instructive to first consider the KK monopole. Since the $d = 10$ KK monopole is T-dual to the type IIB NS5-brane it is a solitonic object whose effective tension scales with $\alpha = -2$. Applying (20) for $\alpha = -2$ and $p = 5$ gives a scaling parameter $\beta = 0$. Since the $d = 10$ KK monopole is related to the $d = 11$ KK monopole via a double dimensional reduction we expect to obtain $\beta = 1$ and not $\beta = 0$. The resolution to this apparent puzzle lies in the fact that the $d = 10$ KK monopole action is not of the standard form (15). Instead it belongs to the more general class of worldvolume actions [12]

$$S(d = 10) = -T \int d^{p+1}x \, e^{\alpha\phi} |k'|^\gamma \sqrt{|\partial_i X^\mu \partial_j X^\nu g_{\mu\nu}|} \quad (21)$$

with $p = 5, \alpha = -2$ and $\gamma = 2$. Here $k'^\mu \neq k^\mu$ is a Killing vector. In the case of the KK monopole the Killing vector k'^μ refers to the Taub–NUT isometry direction. As before, the action (21) can be rewritten, for any value of γ , as a gauged sigma model with an 11-dimensional target space:

$$S(d = 11) = -T \int d^{p+1}\hat{x} \, |\hat{k}|^\beta |\hat{k}'|^\gamma \sqrt{|\partial_i \hat{X}^{\hat{\mu}} \partial_j \hat{X}^{\hat{\nu}} \hat{\Pi}_{\hat{\mu}\hat{\nu}}|}. \quad (22)$$

Equation (20) (corresponding to the special case $\gamma = 0$) gets replaced by the more general formula:

$$\alpha = -\frac{1}{3}(p+1) + \frac{2}{3}(\beta - \frac{1}{2}\gamma). \quad (23)$$

Substituting the appropriate values for p, α and γ corresponding to the $d = 10$ KK-monopole we obtain $\beta = 1$, as expected.

The discussion above is summarized in figure 3. In this figure we have also indicated how the effective tension of the $d = 10$ type IIA KK monopole transforms into the effective tension of the type IIB NS5-brane via T-duality in the Killing isometry direction. Here we have used the standard T-duality rules:

$$|k| \rightarrow \frac{1}{|k|}, \quad e^\phi \rightarrow \frac{1}{|k|} e^\phi. \quad (24)$$

[†] Note that in the case of the D-4-brane/M-5-brane action the BI 1-form must be oxidized to a self-dual 2-form in order not to upset the counting of the bosonic degrees of freedom.

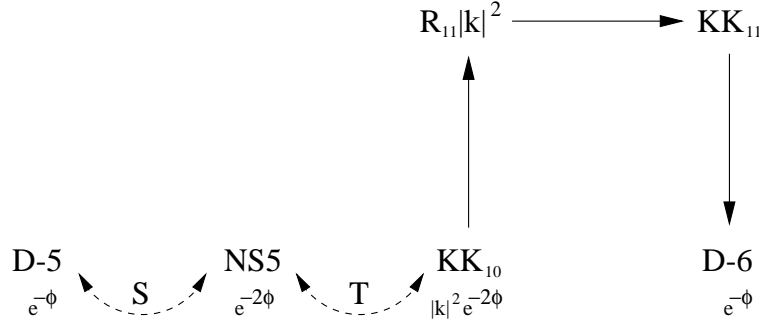


Figure 3. The effective type IIA KK-monopole tension. The arrow pointing up relates the effective action of the IIA KK-monopole and the gauged σ -model action. The arrow to the right relates the gauged σ -model action to the $d = 11$ KK-monopole action and the arrow pointing down represents direct dimensional reduction to the D-6-brane action. The curved arrows indicate (T and S) duality transformations. Below each brane we have indicated how the effective tension scales.

Furthermore we have indicated how the effective tension of the type IIB NS5-brane transforms into the effective tension of the type IIB D-5-brane under a S-duality transformation. In general the effective tension of a D- p -brane (p odd) transforms under S-duality as

$$e^{-\phi} \sqrt{|g|} \rightarrow e^{\delta\phi} \sqrt{|g|} \quad (25)$$

with

$$\delta = -\frac{1}{2}(p-1). \quad (26)$$

The case in figure 3 corresponds to $p = 5$. The exact transformations between the complete worldvolume 5-brane actions is discussed in [21].

We now turn our attention to the M-9-brane. We have seen that the oxidation of the D-8-brane leads to a gauged sigma model that scales as $(R_{11})^3$ (see figure 1). Therefore the M-9-brane cannot move in 11 decompactified dimensions. At this stage we mimic the situation for the KK-monopole and propose the following (Nambu–Goto part of the) worldvolume M-9-brane action:

$$S_{\text{M-9}} = -T \int d^9 \hat{x} \, |\hat{k}|^3 \sqrt{|\partial_i \hat{X}^{\hat{\mu}} \partial_j \hat{X}^{\hat{\nu}} \hat{\Pi}_{\hat{\mu}\hat{\nu}}|}. \quad (27)$$

Note that we do not integrate over the special (compact) isometry direction defined by the Killing vector. Therefore the action (27) describes a *wrapped* 9-brane instead of a 9-brane with 9 non-compact (spacelike) worldvolume directions.

By construction, the reduction of the M-9-brane action (27) over the special isometry direction yields the worldvolume action of the D-8-brane. Similarly, the direct dimensional reduction (i.e. the reduction over the single transversal direction) is expected to yield the worldvolume action of the conjectured (wrapped) IIA-9-brane [7][†]. Applying equation (23), for $(p, \beta, \gamma) = (8, 0, 3)$ [‡] we find $\alpha = -4$. This agrees with the dilaton coupling predicted in [7] by other considerations. We conclude that the M-9-brane action (27) reproduces the correct dilaton couplings for both the type IIA D-8-brane as well as the type IIA 9-brane upon double and direct dimensional reduction, respectively.

[†] Note that at the level of solutions the type IIA 9-brane is represented by $d = 10$ Minkowski spacetime which has unbroken supersymmetry. The presence of an orientifold (see below) breaks half of the supersymmetry.

[‡] Note that β gives the scaling with respect to the 11th direction which in this case is a direction transverse to the brane. Therefore we have $\beta = 0$.

In order to perform the direct dimensional reduction of the M-9-brane we must compactify the single transverse direction. Charge conservation requires that after compactification we include another source of charge such as an orientifold. It would be interesting to see whether, after including this orientifold, the type IIA 9-brane leads to a description of the $E_8 \times E_8$ heterotic superstring theory. This would provide a T-dual version of the description of the $SO(32)$ heterotic superstring via type IIB 9-branes as discussed in [7]. In this way every 10-dimensional 9-brane corresponds to a $N = 1$ superstring theory:

$$\begin{aligned} \text{IIA} - 9 &\leftrightarrow \text{heterotic } E_8 \times E_8, \\ \text{IIB} - 9 &\leftrightarrow \text{heterotic } SO(32), \\ \text{D} - 9 &\leftrightarrow \text{type I } SO(32). \end{aligned}$$

The discussion above is summarized in figure 4. In this figure we have also indicated how the effective tension of the wrapped type IIA 9-brane transforms into the effective tension of the unwrapped type IIB 9-brane, or (1,0)-brane, under a T-duality transformation. Performing a T-duality in the isometry direction of the type IIA 9-brane and applying (24) we first find $|k|^3 e^{-4\phi} \rightarrow |k|^{1/2} e^{-4\phi}$ which is the effective tension of a wrapped type IIB 9-brane. The factor $|k|^{1/2}$ can now be absorbed into the Nambu–Goto term and we obtain an unwrapped type IIB 9-brane. We have also indicated how a S-duality transformation relates the type IIB 9-brane to the D-9-brane, or (0,1)-brane[†]. In the latter S-duality transformation we have used equation (26) for $p = 9$. Note that the T-duality between the wrapped type IIA 9-brane and the unwrapped type IIB 9-brane is similar to the T-duality between the (wrapped) type IIA KK monopole and the unwrapped type IIB NS5-brane. In both cases one brane has a special Killing vector (type IIA KK monopole, type IIA 9-brane) whereas the T-dual brane has no such Killing vector (type IIB NS5-brane, type IIB 9-brane). Figure 4 shows that S- and T-duality and the existence of the D-9-brane provide another argument in favour of the type IIA 9-brane and its M-theory origin, the M-9-brane.

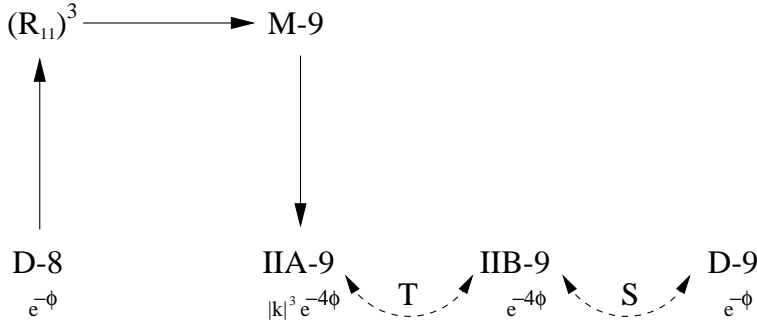


Figure 4. The M-9-brane effective tension. The arrow pointing up relates the effective action of the D-8-brane and the gauged σ -model action. The arrow to the right relates the gauged σ -model action to the M-9-brane action and the arrow pointing down represents direct dimensional reduction to the (wrapped) type IIA 9-brane action. The curved arrows indicate (T and S) duality transformations. Below each brane we have indicated how the effective tension scales.

So far, we have only discussed the leading Nambu–Goto term of the M-9-brane worldvolume action. In principle one should be able to reconstruct the complete action by using the relation, via reduction and duality, with the worldvolume actions of the other branes given in figure 4. Since we are dealing with a wrapped M-9-brane we have a nine-dimensional worldvolume and therefore we expect the worldvolume fields to form a $d = 9$

[†] The (1,0)-brane and (0,1)-brane are two special cases of a whole family of (p, q) 9-branes [7].

vector multiplet. The single scalar of this vector multiplet describes the dynamics in the single transverse direction in which the M-9-brane can move. It would be interesting to explicitly construct the complete M-9-brane worldvolume action.

3. The D-8-brane target space solution

Before discussing the M-9-brane target space solution in the next section we will first review in this section some properties of the D-8-brane target space solution. The D-8-brane target space solution arises as a domain wall solution[†] in massive IIA supergravity [9]. When the cosmological constant term is dualized to a 9-form potential one can write down a domain wall solution with different values of the cosmological constant at different sides of the domain wall [4]. This solution is coordinate equivalent to the conformally flat solution given in [3].

We consider the following ansatz for the extreme D-8-brane solution (in string frame metric):

$$\begin{aligned} ds_{10}^2 &= H^\alpha (dt^2 - dx_{(8)}^2) - H^\beta dy^2, \\ e^{2\phi} &= H^\gamma \quad C_{012345678}^{(9)} = H^\epsilon, \end{aligned} \quad (28)$$

where $H = H(y)$ is a harmonic function over the single transverse direction y whose form, in a local neighbourhood, is given by

$$H(y) = c + Q|y| \quad (29)$$

in terms of two constants c and Q . In order to avoid a singularity at $H = 0$ and to obtain a real dilaton, we use the absolute value of y in the harmonic function and take $c > 0$ and $Q > 0$. The ansatz (28) describes a domain wall positioned at $y = 0$.

It turns out that the parameter ϵ cannot be determined by the equations of motion obtained by minimizing the action (12). This is in contradistinction to the D- p -branes with $p < 8$ where one finds, for all $p < 8$, that $\epsilon = -1$. Solving the equations of motion obtained from the action (12) we find the following expressions for α, β, γ :

$$\alpha = \frac{1}{2}\epsilon, \quad \beta = -\frac{5}{2}\epsilon - 2, \quad \gamma = \frac{5}{2}\epsilon. \quad (30)$$

Furthermore, substituting the solution (30) into (14), we find the following relation between m and Q :

$$m = \pm\epsilon Q. \quad (31)$$

Notice that for any fixed m and Q we find two solutions corresponding to ϵ and $-\epsilon$. For $m = 0$ there is a single solution with $\epsilon = 0$ or, equivalently, $Q = 0$. In this case the metric reduces to that of a flat Minkowski spacetime.

The value of the cosmological constant differs at the two sides of the domain wall whenever we make different choices for the constant Q at the left and right of the domain wall:

$$\begin{aligned} H(y) &= c + Q_L|y|, & y < 0, \\ H(y) &= c + Q_R|y|, & y > 0. \end{aligned} \quad (32)$$

The free parameter ϵ labelling the above D-8-brane solutions is related to the freedom to perform a coordinate transformation in y keeping the solution within the ansatz (28). To show this we first note that at one side of the domain wall one can always use coordinates such that $y \geq 0$. By performing a suitable shift transformation, $y' = y + c/Q$, the harmonic function can always be written as $H(y') = Qy'$, where $y' \in (c/Q, \infty)$ and $Q > 0$. We only consider coordinate transformations that keep the transversal coordinate y within this (positive

[†] For earlier work on domain wall solutions in $d = 4$ supergravity theories, see [22].

and infinite) range. Consider a D-8-brane for a given negative value of ϵ , say $\epsilon = -1$ or $m = \pm Q$. We perform the following coordinate transformation labelled by ϵ (we omit the prime on the shifted coordinate):

$$y \rightarrow y' = f(\epsilon)y^{-\frac{1}{\epsilon}}, \quad (33)$$

with the function $f(\epsilon)$ given by

$$f(\epsilon) = -\epsilon Q^{-\frac{1+\epsilon}{\epsilon}}. \quad (34)$$

We restrict ourselves to negative ϵ (for positive ϵ the range of y' would become negative and finite). Under the above coordinate transformation the harmonic function H transforms as

$$H_{(\epsilon=-1)}(y) = -my = (H_{(\epsilon)}(y'))^{-\epsilon} = (Q'y')^{-\epsilon}, \quad (35)$$

with $Q' = -Q/\epsilon$. Therefore all solutions with negative ϵ , defined at positive y , are related to each other by the coordinate transformations (33). The same holds for all positive ϵ solutions. Starting with a solution for a given positive value of ϵ , say $\epsilon = 1$, we obtain all other positive ϵ solutions by performing the transformation (33) with ϵ replaced by $-\epsilon$.

Finally, to relate solutions with positive and negative values of ϵ one must perform a coordinate transformation of the form

$$y \rightarrow 1/y. \quad (36)$$

This transformation is included in (33) if we also allow positive values of ϵ . However, under the transformation (36) the infinite positive domain of y gets mapped to a finite negative domain in y' and Q becomes negative to keep $H(y)$ positive. This means that in the new coordinate system $y = 0$ acts as infinity. In any case, we see that for all values of ϵ , positive and negative, the domain wall solutions are coordinate equivalent to each other. In the remaining part of this work we will often concentrate on the $\epsilon < 0$ solutions because in that case the position of the domain wall occurs for a finite value of y and the asymptotic region far away from the domain wall is reached by taking $y \rightarrow \infty$.

We next consider the asymptotic values of the D-8-brane configuration, keeping in mind that solutions labelled by different values of ϵ are coordinate equivalent. For reasons explained in the previous paragraph we restrict ourselves to $\epsilon < 0$. We first consider the dilaton. For $\epsilon < 0$, the string coupling e^ϕ approaches zero in the limit $y \rightarrow \infty^\dagger$ and $c^{5\epsilon/4}$ in the limit $y \rightarrow 0$ (the position of the domain wall). Because of the freedom in choosing c the string coupling can take any value at the position of the D-8-brane, when $y \rightarrow 0$. In particular, it becomes infinite if $c \rightarrow 0$.

To determine the asymptotic structure of the metric we calculate the Riemann tensor squared. For general (positive and negative) ϵ one finds [23]:

$$R_{\mu\nu\delta\rho}R^{\mu\nu\delta\rho} \propto H^{5\epsilon}(\partial_y H)^4. \quad (37)$$

This shows that, assuming that $\epsilon < 0$, in the limit $y \rightarrow \infty$ the curvature approaches zero, while the curvature near the D-8-brane position approaches a constant as long as we keep $c > 0$. If we take the limit $c \rightarrow 0$, the curvature blows up and we are dealing with a true singularity.

It is instructive to consider timelike geodesics in the background of a D-8-brane. The equation of such a geodesic is given by (in Einstein frame):

$$\ddot{y} = -\frac{1}{16}\epsilon Q H(y)^{3\epsilon+1}. \quad (38)$$

We deduce that for $\epsilon < 0$ test-objects will be repelled by the D-8-brane. Only when $\epsilon < -1/3$ will the acceleration approach zero when $y \rightarrow \infty$. We can solve (38) for general ϵ and find

[†] This is different from the other D- p -branes ($0 \leq p \leq 6$) where the string coupling approaches a constant at transverse infinity.

that, when a test-particle is inserted in a D-8-brane geometry with zero velocity, the geodesic trajectory is described by

$$y(t) = (at^2 + b)^{-1/3\epsilon} - c, \quad (39)$$

where $a > 0$ and b are constants determined by $y(0)$, Q and m^2 and c is the constant in the harmonic function $H(y) = c + Q|y|$. This formula shows the difference between the $\epsilon > 0$ and $\epsilon < 0$ solutions. For positive ϵ the geodesics approach a finite value of y if we take t to infinity. For negative ϵ the test-particles move towards infinity if we take t to infinity.

Finally, for special values of $\epsilon < 0$ the D-8-brane solution reduces to a more familiar form. We mention the following three choices:

(a) $\epsilon = -1$ gives

$$\begin{aligned} ds_{10}^2 &= H^{-1/2}(dt^2 - dx_{(8)}^2) - H^{1/2}dy^2, \\ e^{2\phi} &= H^{-5/2}, \quad C_{012345678}^{(9)} = H^{-1}. \end{aligned} \quad (40)$$

This is the standard form of the D-8-brane solution as it is obtained via T-duality from the other D- p -brane solutions.

(b) $\epsilon = -\frac{2}{3}$ gives

$$\begin{aligned} ds_{10}^2 &= H^{-1/3}(dt^2 - dx_{(8)}^2 - dy^2), \\ e^{2\phi} &= H^{-5/3}, \quad C_{012345678}^{(9)} = H^{-2/3}. \end{aligned} \quad (41)$$

This is the conformal flat metric solution given in [3].

(c) $\epsilon = -\frac{4}{5}$ gives

$$\begin{aligned} ds_{10}^2 &= H^{-2/5}(dt^2 - dx_{(8)}^2) - dy^2, \\ e^{2\phi} &= H^{-2}, \quad C_{012345678}^{(9)} = H^{-4/5}. \end{aligned} \quad (42)$$

The special feature of this choice of ϵ is that $\beta = 0$.

4. The M-9-brane target space solution

In order to oxidize the D-8-brane to an M-theory solution we must add a cosmological constant, proportional to m^2 , to the $d = 11$ supergravity theory. As discussed in the introduction, to do this we must assume that the $d = 11$ background fields have an isometry generated by a Killing vector k^μ (in this section we omit hats on the fields). The extra 11th direction corresponds to this isometry direction. Notice that the present situation differs from the oxidation of the D-6-brane to an M-theory Kaluza–Klein monopole. In the case of the monopole the isometry direction in the Taub–NUT space is required in order to solve the equations of motion of massless $d = 11$ supergravity. Here, the extra isometry direction is already needed, before solving the equations of motion, in order to write down the massive $d = 11$ supergravity Lagrangian.

Although a $d = 11$ manifold with an isometry is basically a $d = 10$ manifold, the $d = 10$ interpretation of a solution may differ from the $d = 11$ interpretation of the same solution. An example of this is the $d = 11$ KK monopole. Its $d = 10$ interpretation is the D-6-brane which is singular. The $d = 11$ interpretation of the same solution leads to a non-singular solution. In order to show the regularity of the solution one crucially needs the compact isometry direction. The singularity of the 10-dimensional D-6-brane can be viewed as an illegitimate neglect of KK modes which become massless at the D-6-brane core (see e.g. [6]). In the case of the M-9-brane the situation is different since, in order to write down an 11-dimensional action, we have already restricted ourselves to background fields with an isometry.

Assuming the extra isometry direction it is straightforward to rewrite the D-8-brane metric and dilaton in terms of the M-theory metric. We use the following reduction ansatz for the $d = 11$ metric:

$$ds_{11}^2 = e^{-\frac{2}{3}\phi} ds_{10}^2 + e^{\frac{4}{3}\phi} dz^2, \quad (43)$$

where z is the isometry direction and ds_{10}^2 is the 10-dimensional string frame metric. Applying this formula we obtain the following oxidized D-8-brane, or (wrapped) M-9-brane, solution:

$$ds_{11}^2 = H^{-\frac{1}{3}\epsilon} (dt^2 - dx_{(8)}^2) - H^{-\frac{10}{3}\epsilon-2} dy^2 - H^{\frac{5}{3}\epsilon} dz^2, \quad (44)$$

$$H(y) = c + Q|y| \quad (c, Q > 0), \quad m = \pm\epsilon Q.$$

This M-9-brane solution, when reduced over the z -direction, gives the D-8-brane solution of the previous section. The arbitrary parameter ϵ is related to a coordinate transformation in y , as explained in the previous section.

For generic values of ϵ the M-9-brane metric given in (44) represents a 3-block solution, i.e. it exhibits three inequivalent directions. This is different from the M-2-brane and M-5-brane which are represented by 2-block solutions with the two inequivalent directions corresponding to the worldvolume and transverse directions. Only for special values of ϵ will the M-9-brane metric also reduce to a 2-block solution:

(a) $\epsilon = -\frac{2}{5}$ gives

$$ds_{11}^2 = H^{2/15} (dt^2 - dx_{(8)}^2) - H^{-2/3} (dy^2 + dz^2). \quad (45)$$

(b) $\epsilon = -\frac{2}{3}$ gives

$$ds_{11}^2 = H^{2/9} (dt^2 - dx_{(8)}^2 - dy^2) - H^{-10/9} dz^2. \quad (46)$$

The first case gives a metric suggesting z represents a special isometry direction transverse to the brane. This is what happens for the M-theory Kaluza–Klein monopole.

We first consider in more detail the properties of a single M-9-brane. Next, we will study a system of two M-9-branes.

4.1. The single M-9-brane

We consider a single domain wall solution, positioned at $y = 0$, with harmonic function given by (32). Since $g_{zz} = (R_{11})^2$ it follows from the M-9-brane metric that $R_{11} = H(y)^{\frac{5}{6}\epsilon}$. We take $\epsilon < 0$ for reasons explained in the previous section. We conclude that for an observer far away from the M-9-brane, spacetime is effectively 10-dimensional: $R_{11} \rightarrow 0$ if $y \rightarrow \infty$. On the other hand, close to the M-9-brane position R_{11} approaches a constant value: $R_{11} \rightarrow c^{\frac{5}{6}\epsilon}$ if $y \rightarrow 0$. For a fixed (negative) value of ϵ the radius at the position of the domain wall becomes very large if the value of c is chosen to be small: $R_{11} \rightarrow \infty$ if $c \rightarrow 0^+$. Thus we see that for a general M-9-brane the asymptotic geometry ($|y| \rightarrow \infty$) looks 10-dimensional and only when we approach the domain wall does the 11th dimension open up.

We now consider a special case in which the cosmological constant is zero at one side of the M-9-brane (say for $y < 0$), i.e.

$$H(y) = c \quad y < 0, \quad (47)$$

$$H(y) = c + Q_R|y| \quad y > 0.$$

† Note that the curvature at $y = 0$ becomes infinite in the limit $c \rightarrow 0$.

d=11

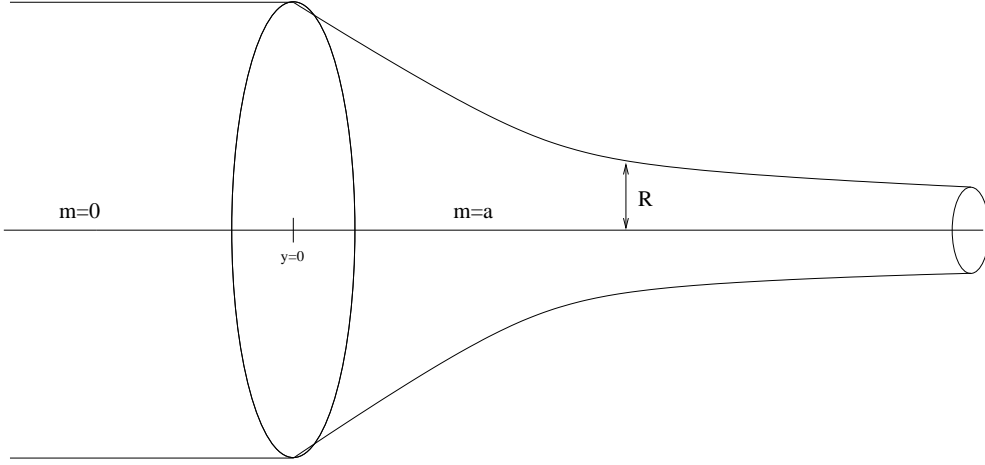


Figure 5. The single M-9-brane. The cosmological constant has been taken equal to zero at the left of the domain wall and non-zero at the right. For $c \neq 0$ ($c \rightarrow 0$) the radius of the 11th dimension R_{11} is finite (infinite) at $y = 0$. In the limit $c \rightarrow 0$ the asymptotic spacetime is given by 10-dimensional (11-dimensional) Minkowski space for $y \rightarrow \infty$ ($y \rightarrow -\infty$).

In that case the $d = 11$ M-theory metric is given by

$$\begin{aligned} ds_{11}^2 &= c^{-\frac{1}{3}\epsilon} (dt^2 - dx_{(8)}^2) - c^{-\frac{10}{3}\epsilon-2} dy^2 - c^{\frac{5}{3}\epsilon} dz^2, & y < 0, \\ ds_{11}^2 &= H^{-\frac{1}{3}\epsilon} (dt^2 - dx_{(8)}^2) - H^{-\frac{10}{3}\epsilon-2} dy^2 - H^{\frac{5}{3}\epsilon} dz^2, & y > 0. \end{aligned} \quad (48)$$

Thus, we see that at the $y < 0$ side of the domain wall, the $d = 11$ spacetime is given by the direct product of a $d = 10$ Minkowski spacetime and a circle with radius $R_{11} = c^{5\epsilon/6}$. Since $\epsilon < 0$ we obtain in the limit $c \rightarrow 0$ an *unwrapped* domain wall whose geometry at one side is given by a decompactified $d = 11$ Minkowski spacetime and at the other side by a $d = 11$ geometry with an isometry which, far away from the domain wall, looks like a 10-dimensional spacetime (see figure 5). The domain wall interpolates between two spacetimes which are related to each other by a decompactification of the special z direction.

4.2. Two M-9-branes

We next consider two parallel M-9-branes where each of the two domain walls have the product space $M^{10} \times S^1$ at one side, as described above. To obtain a static solution the two M-9-branes must have the same ‘charge’. In order for this to be the case the $M^{10} \times S^1$ spacetime must lie in between the two M-9-branes and the cosmological constants at the other side of the two M-9-branes must be equal to each other (see figure 6). To be precise

$$\begin{aligned} H_1(y) &= Q|y - y_1| + c, & H_2(y) &= c, & y < y_1, \\ H_1(y) &= H_2(y) = c, & & & y_1 < y < y_2, \\ H_1(y) &= c, & H_2(y) &= Q|y - y_2|, & y > y_2, \end{aligned} \quad (49)$$

where $y = y_1$ (y_2) is the position of the first (second) domain wall.

Consider the region in between the two M-9-branes. In this region we can take the limit $R_{11} \rightarrow \infty$ because $m = 0$. This allows us to *unwrap* the two M-9-branes and to define a massless bulk supergravity theory in between the two M-9-branes. In this way we obtain a system that is reminiscent to the 9-branes of Hořava and Witten [8].

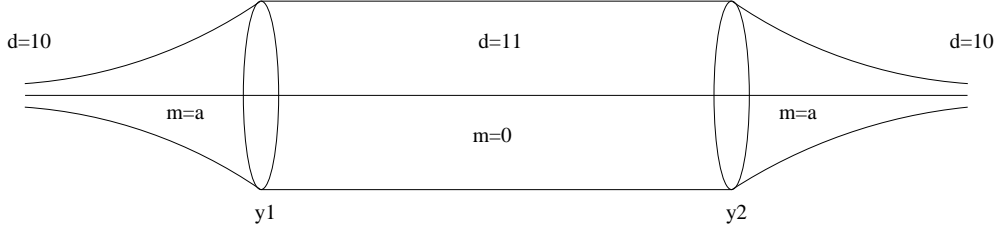


Figure 6. Two M-9-branes. The radius of the 11th dimension R_{11} is finite in between the two domain walls. In this region spacetime can be decompactified in the z direction by taking the limit $c \rightarrow 0$ or $R_{11} \rightarrow \infty$.

In [8] 11-dimensional supergravity is considered on the manifold with boundary $\mathbb{R}^{10} \times \mathbb{S}^1/\mathbb{Z}_2 = \mathbb{R}^{10} \times I$ with I the unit interval. The boundary of this manifold consists of two ‘end of the world’ 9-branes. Anomaly considerations determine that the worldvolume theory of each 9-brane is given by a $N = 1$ vector multiplet with gauge group E_8 . The bulk theory in between the two 9-branes is given by massless $d = 11$ supergravity. According to [8] the strong coupling limit of the $E_8 \times E_8$ heterotic string is given by massless $d = 11$ supergravity on $\mathbb{R}^{10} \times I$ in the same way as the strong coupling behaviour of the type IIA superstring is given by massless $d = 11$ supergravity on $\mathbb{R}^{10} \times \mathbb{S}^1$. In this approach the heterotic dilaton is related to the radius R_{11} in the y -direction[†]. In the limit $R_{11} \rightarrow 0$ the two 9-branes coincide and we end up with a $d = 10$ spacetime.

In the conclusions section we will discuss the similarities and differences between the two-M-9-brane system and the Hořava–Witten system sketched above.

5. Supersymmetry

We first consider the supersymmetry preserved by the D-8-brane. The relevant part of the supersymmetry rules (with parameter η) of massive IIA supergravity is given by (in string frame) [26]

$$\begin{aligned} \delta\psi_\mu &= \partial_\mu\eta - \frac{1}{4}\omega_\mu{}^{ab}\Gamma_{ab}\eta + \frac{1}{8}me^\phi\Gamma_\star\Gamma_\mu\eta, \\ \delta\lambda &= \Gamma^\mu(\partial_\mu\phi)\eta - \frac{5}{4}me^\phi\Gamma_\star\eta. \end{aligned} \quad (50)$$

Substituting the D-8-brane solution (28) into the supersymmetry-preserving conditions $\delta\psi_\mu = \delta\lambda = 0$, we find that the D-8-brane, for each value of ϵ , breaks half of the bulk supersymmetry, with the Killing spinor given by [3, 4]

$$(1 + \Gamma_{012345678})\eta = 0, \quad \eta = H^{\epsilon/8}\eta_0 \quad (51)$$

for constant η_0 .

Substituting the D-8-brane solution into the supersymmetry transformation rules (50) we find an overall factor $H(y)^{\frac{5\epsilon}{4}}$. For negative ϵ we see that in the limit $y \rightarrow \infty$ we have unbroken supersymmetry. We do not find supersymmetry enhancement in the other asymptotic region, i.e. at the position of the domain wall.

[†] Note that in the double dimensional reduction leading to the D-8-brane the IIA dilaton is related to the radius R_{11} in the z -direction.

We next consider the supersymmetry preserved by the M-9-brane. The $d = 10$ massive supersymmetry rules given in equation (50) correspond to the following supersymmetry rule of the $d = 11$ gravitino (we omit hats):

$$\delta\psi_\mu = \partial_\mu\eta - \frac{1}{4}\omega_\mu{}^{ab}\Gamma_{ab}\eta + \frac{i}{24}m|k|(2k^\nu\Gamma_\nu\Gamma_\mu - 3k_\mu)\eta. \quad (52)$$

We substitute the M-9-brane solution (44) into the supersymmetry preserving condition $\delta\psi_\mu = 0$ and find the conditions

$$(1 + i\Gamma_z\Gamma_y)\eta = 0, \quad \eta = H^{-\epsilon/12}\eta_0 \quad (53)$$

for constant η_0 . Assuming that we reduce over the z -direction we have $\Gamma_z = i\Gamma_\star$. The projection operator $i\Gamma_\star\Gamma_y$ is equivalent to the projection operator $\Gamma_{01\dots 8}$ which is not yet of the form (2). However, after performing the redefinition $\Gamma_\mu \rightarrow \Gamma'_\mu = i\Gamma_\star\Gamma_\mu^\dagger$, we obtain

$$(1 + \Gamma_{01\dots 8z})\eta = 0, \quad (54)$$

which is exactly the projection operator of an M-9-brane as suggested by the M-theory supersymmetry algebra; see equation (2).

6. Conclusions

In this work we have investigated some properties of the M-9-brane whose existence is suggested by the structure of the M-theory superalgebra [6, 7]. We have studied the M-9-brane from the point of view of both the worldvolume action as well as the target space solution. We have shown how the action and solution are related, via reduction and duality, to the actions and solutions of $d = 10$ superstring theory.

We expect that the complete M-9-brane action can be constructed by making use of the connection with the known D-9-brane action. Concerning the target space solutions we note that there are several similarities but also differences between the two M-9-brane system considered in this work and the two 9-brane system of Hořava–Witten [8]. One similarity is that the $d = 10$ worldvolume theory of both the M-9-brane and the 9-brane of [8] have $N = 1$ chiral supersymmetry. Another suggestive similarity is that intersections of an M-9-brane with one of the other branes of M-theory (see figure 1) leads to only 1-branes and 5-branes which are exactly the basic objects in $E_8 \times E_8$ heterotic superstring theory [24]. The possible intersections are [25]:

$$(1|W, M9), \quad (1|M2, M9), \quad (5|M5, M9), \quad (5|KK, M9). \quad (55)$$

A difference between our two M-9-brane system and the one of [8] is that in our case the two-domain-wall system arises as a solution to the equations of motion of massive $d = 11$ supergravity whereas in [8] the same configuration is taken as a fixed background for the massless supergravity theory. In our case the $d = 11$ spacetime extends behind the domain walls (see figure 6), whereas the 9-branes of [8] are ‘one-sided’ 9-branes positioned at the end of the world † .

Another, related, difference is that in our case, as we have shown in Section 5, the $d = 11$ bulk supersymmetry is broken by the M-9-brane solution. On the other hand, in the two-9-brane system of [8] supersymmetry is broken by the boundary conditions on the 11-dimensional fields.

It remains to be further investigated whether or not the two-M-9-brane system we have considered in this work is related to the two ‘end of the world’ 9-branes of Hořava and Witten. We hope to report on progress in this direction in the near future.

† Using the Killing vector this redefinition can also be written in an 11-dimensional notation.

‡ Recently, one-sided domain wall solutions have been investigated (G Gibbons, private communication).

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